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Effective Kähler Potentials

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Abstract

We compute the 1-loop effective Kähler potential in the most general renormalizable $N = 1$ $d = 4$ supersymmetric quantum field theory.

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1 Introduction

The effective action of a theory incorporates all its quantum corrections. It is, however, nonlocal and difficult to calculate even in exactly solvable models. Consequently, various approximations are used to evaluate parts of the effective action. In particular, one frequently uses a loop expansion and a momentum expansion. The leading term in the momentum expansion is the effective potential [1], and it determines the vacuum structure of the theory. The next term is of order p^2 , and, together with the effective potential, determines the masses of the states in the theory.

In $N = 1$ supersymmetric theories, there is a nonrenormalization theorem that protects the superpotential from corrections; the calculation that is directly analogous to calculations of the effective potential is the calculation of the effective Kähler potential [2, 3, 4]. When supersymmetry is unbroken, this quantity determines both the effective potential as well the normalization of the p^2 term. In this paper, we compute the 1-loop effective Kähler potential in the most general renormalizable $N = 1$ $d = 4$ supersymmetric quantum field theory. Our methods are considerably simpler than earlier analyses [2, 3], and our results are more general and complete; they take a particularly nice form in supersymmetric Landau gauge. We find this encouraging, since it has been argued that, at least to one-loop, the Landau gauge effective action is equivalent to the DeWitt-Vilkovisky gauge independent effective action [5, 6, 7].

The paper is organized as follows: After describing the most general renormalizable $N = 1$ $d = 4$ supersymmetric theory, we compute the 1-loop effective Kähler potential using standard super-Feynman rules in super-Landau gauge. We then use functional methods to generalize our result to arbitrary values of the gauge parameter. Finally, we study some $N = 2$ examples: the contribution of a massive hypermultiplet to the vector multiplet low energy effective action, and the effective Kähler potential for the hypermultiplet itself.

2 Setup

We start from the most general renormalizable $N = 1$ $d = 4$ supersymmetric action (we use the conventions of [8]):

$$S = \int d^4x d^2\theta \frac{1}{8g^2} W_A^\alpha W_\alpha^A + \int d^4x d^4\theta \bar{\Phi} e^V \Phi + \left(\int d^4x d^2\theta \left(\frac{1}{2} \Phi^i m_{ik} \Phi^k + \frac{1}{3!} \lambda_{ijk} \Phi^i \Phi^j \Phi^k \right) + h.c. \right). \quad (2.1)$$

In this expression the chiral superfields Φ are in some product representation of a big gauge group consisting of all the relevant gauge groups in the problem. For example, in the $N = 2$ case with a matter hypermultiplet Q, \tilde{Q} in a representation R of the gauge group and the chiral component superfield ϕ of the vector multiplet, Φ would be in a direct sum of the adjoint representation with R and the conjugate representation \tilde{R} :

$$\Phi = \begin{pmatrix} \phi_A \\ Q_i \\ \tilde{Q}^k \end{pmatrix}. \quad (2.2)$$

The mass matrix m as well as the couplings λ are gauge invariant and symmetric. Again, in the $N = 2$ case, m would look like

$$m = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & m_H \delta_{ik} \\ 0 & m_H \delta_{ki} & 0 \end{pmatrix}, \quad (2.3)$$

where m_H is the mass of the hypermultiplet and λ would be such that $\phi^A \lambda_{Aik} = \phi^A (T_A^R)_{ik}$ and $\bar{\phi}^A \lambda_{Aik} = \bar{\phi}^A (T_A^{\tilde{R}})_{ik}$

3 Feynman graph calculation

We now give an explicit Feynman graph calculation of the 1-loop effective Kähler potential. As we shall see, this is simplest in supersymmetric Landau gauge.

To calculate the one loop Kähler potential, we expand the action to quadratic order around some constant (x, θ -independent) background value

Φ of the scalar fields, *i.e.*, $\Phi = \bar{\Phi} + \varphi$. Propagators are defined from kinetic terms that are independent of the background value $\bar{\Phi}$ of Φ , and vertices are defined from the remaining mass and $\bar{\Phi}$ dependent terms. It is convenient to define a $\bar{\Phi}$ dependent $\varphi\varphi$ mass μ that contains all the m and λ dependence, as well as $\bar{\Phi}$ dependent φV and VV masses X and M :

$$\begin{aligned}\mu_{ik} &= m_{ik} + \lambda_{ijk} \bar{\Phi}^j \\ X_i^A &= 2g (T_A)_i^k \bar{\Phi}_k \\ \bar{X}_A^i &= 2g \bar{\Phi}^k (T_A)_k^i \\ M_{AB} &= \frac{1}{2} \left(\bar{X}_A^i X_{Bi} + \bar{X}_B^i X_{Ai} \right),\end{aligned}\tag{3.1}$$

where we have rescaled V by a factor $2g$. Adding a supersymmetric gauge fixing term $\xi^{-1} D^2 V \bar{D}^2 V$, we can write the general gauge action as

$$\begin{aligned}\frac{1}{2} \int d^4x d^4\theta \left(V^A \left(D^\alpha \bar{D}^2 D_\alpha - \frac{1}{\xi} \{ D^2, \bar{D}^2 \} \right) V_A + 2\bar{\varphi}^i \varphi_i + 2V^A \bar{X}_A^i \varphi_i \right. \\ \left. + 2\bar{\varphi}^i X_i^A V_A + V^A M_{AB} V^B \right) + \frac{1}{2} \left(\int d^4x d^2\theta \varphi^i \mu_{ik} \varphi^k + h.c. \right).\end{aligned}\tag{3.2}$$

In principle there are also ghosts in the action, but since they do not couple to the scalar fields, to one-loop we need not consider those terms. The effective Kähler potential is most conveniently calculated in supersymmetric Landau gauge $\xi = 0$. In this gauge the VV propagator becomes $-D^\alpha \bar{D}^2 D_\alpha / \square^2$, which implies that by D -algebra at one-loop there can be no mixed contributions, *i.e.*, loops containing both VV and $\varphi\bar{\varphi}$ propagators. This leaves only two basic loops that have to be calculated. One is the sum over diagrams with n external μ vertices and n external $\bar{\mu}$ vertices and $2n$ scalar field propagators $-1/\square$. After Fourier transforming and summing up all the diagrams we get

$$\int \frac{d^4k d^4\theta}{(2\pi)^4} \frac{1}{2k^2} \text{Tr} \ln \left(1 + \frac{\bar{\mu}\mu}{k^2} \right) = \frac{1}{2(4\pi)^2} \text{Tr} \int d^4\theta dk^2 \ln \left(1 + \frac{\bar{\mu}\mu}{k^2} \right).\tag{3.3}$$

The other is the sum over diagrams with n external M vertices and n vector propagators $-D^\alpha \bar{D}^2 D_\alpha / \square$ which, after Fourier transforming and doing the

sum, gives

$$- \int \frac{d^4 k d^4 \theta}{(2\pi)^4} \frac{1}{k^2} \text{Tr} \ln \left(1 + \frac{M}{k^2} \right) = - \frac{1}{(4\pi)^2} \text{Tr} \int d^4 \theta dk^2 \ln \left(1 + \frac{M}{k^2} \right) . \quad (3.4)$$

Evaluating the momentum integral with an ultraviolet cut-off Λ , we find the regularized 1-loop effective Kähler potential for the most general renormalizable four-dimensional theory:

$$K_{eff} = - \frac{1}{2(4\pi)^2} \text{Tr} \left[\bar{\mu} \mu \ln \left(\frac{\bar{\mu} \mu}{e \Lambda^2} \right) - 2M \ln \left(\frac{M}{e \Lambda^2} \right) \right] . \quad (3.5)$$

Note that the factor of $e = \exp(1)$ can be removed by a finite renormalization.

4 General gauge

The effective action, being an off-shell quantity, is not gauge independent. For completeness, we now compute the 1-loop effective Kähler potential for arbitrary gauge fixing parameter ξ using functional methods. To functionally integrate over the chiral fields, it is convenient to solve the chirality constraint on φ and $\bar{\varphi}$ by introducing unconstrained fields ψ and $\bar{\psi}$ such that $\varphi = \bar{D}^2 \psi$ and $\bar{\varphi} = D^2 \bar{\psi}$. In principle this introduces a new gauge invariance into the action, but in the absence of background gauge fields V , the ghosts associated with this gauge fixing decouple (actually, covariant functional methods can also be used in a V background, but they are not needed here [8, 9]). The unconstrained functional integration over V, ψ and $\bar{\psi}$ gives a contribution to the effective action:

$$- \frac{1}{2} \int d^4 x d^4 \theta d^4 \theta' \delta(\theta' - \theta) \text{Tr} \ln \mathcal{O}(x, \theta) \delta(\theta - \theta') , \quad (4.1)$$

where the $-1/2$ comes from the fact that we are computing $\det^{-\frac{1}{2}}$ and \mathcal{O} is the kinetic operator

$$\mathcal{O} = \begin{pmatrix} D^\alpha \bar{D}^2 D_\alpha - \frac{1}{\xi} \{D^2, \bar{D}^2\} + M & \bar{X} \bar{D}^2 & X D^2 \\ & \bar{X} \bar{D}^2 & \mu \bar{D}^2 & \square \\ & X D^2 & \square & \bar{\mu} D^2 \end{pmatrix} . \quad (4.2)$$

The logarithm can be split into two pieces using the formula for the determinant of a block matrix

$$\det \begin{pmatrix} A & B \\ C & E \end{pmatrix} = \det(E) \det(A - B E^{-1} C) . \quad (4.3)$$

To do this we need the formula for the inverse of E which in our case is

$$E^{-1} = \begin{pmatrix} \mu \bar{D}^2 & \square \\ \square & \bar{\mu} D^2 \end{pmatrix}^{-1} = \frac{1}{\square} \begin{pmatrix} -\bar{\mu} \frac{D^2}{\square - \mu \bar{\mu}} & 1 + \bar{\mu} \mu \frac{D^2 \bar{D}^2}{\square(\square - \mu \bar{\mu})} \\ 1 + \mu \bar{\mu} \frac{\bar{D}^2 D^2}{\square(\square - \mu \bar{\mu})} & -\mu \frac{\bar{D}^2}{\square - \mu \bar{\mu}} \end{pmatrix} . \quad (4.4)$$

Using this, we can write

$$\begin{aligned} \text{Tr} \ln(\mathcal{O}) &= \text{Tr} \ln(E) \\ &+ \text{Tr} \ln \left(D^\alpha \bar{D}^2 D_\alpha - \frac{1}{\xi} \{D^2, \bar{D}^2\} + M_{AB} \right. \\ &\quad - \frac{\bar{D}^2 D^2}{\square} \bar{X}_A^i X_{Bi} - \frac{D^2 \bar{D}^2}{\square} X_A^i \bar{X}_{Bi} \\ &\quad - \frac{\bar{D}^2 D^2}{\square} \bar{X}_A^i \left[\bar{\mu} \mu \frac{1}{\square - \mu \bar{\mu}} \right]_{ik} X_B^k \\ &\quad - \frac{D^2 \bar{D}^2}{\square} X_A^i \left[\mu \bar{\mu} \frac{1}{\square - \mu \bar{\mu}} \right]_{ik} \bar{X}_B^k \\ &\quad \left. + \bar{D}^2 \bar{X}_A^i \left[\bar{\mu} \frac{1}{\square - \mu \bar{\mu}} \right]_{ik} \bar{X}_B^k + D^2 X_A^i \left[\mu \frac{1}{\square - \mu \bar{\mu}} \right]_{ik} X_B^k \right) . \end{aligned} \quad (4.5)$$

Since we are only interested in the Φ dependence, we can factor out the Φ independent V -propagator piece and subsequently drop it. Doing this and using the symmetry of μ and $\bar{\mu}$ we have

$$\begin{aligned} \text{Tr} \ln(\mathcal{O}) &= \text{Tr} \ln(E) \\ &+ \text{Tr} \ln \left(1 + \frac{D^\alpha \bar{D}^2 D_\alpha}{\square^2} M_{AB} + \xi \frac{\bar{D}^2 D^2}{\square^2} (S_{AB} + T_{AB}) \right) \end{aligned}$$

$$+ \xi \frac{D^2 \bar{D}^2}{\square^2} (S_{BA} + T_{BA}) - \xi \frac{D^2}{\square} R_{AB} - \xi \frac{\bar{D}^2}{\square} \bar{R}_{AB} \Big) , \quad (4.6)$$

where

$$\begin{aligned} S_{AB} &= \frac{1}{2} \left(\bar{X}_A^i X_{Bi} - \bar{X}_B^i X_{Ai} \right) , \\ T_{AB} &= \bar{X}_A^i \left[\bar{\mu} \mu \frac{1}{\square - \bar{\mu} \mu} \right]_{ik} X_B^k , \\ R_{AB} &= \frac{1}{2} X_{(A}^i \left[\mu \frac{1}{\square - \bar{\mu} \mu} \right]_{ik} X_{B)}^k . \end{aligned} \quad (4.7)$$

Because the operator $D^\alpha \bar{D}^2 D_\alpha$ annihilates anything proportional to D^2 or \bar{D}^2 , this second trace actually splits up into a sum of two terms, and we can thus write (4.5) as a sum of three terms

$$\begin{aligned} \text{Tr} \ln \mathcal{O} &= \text{Tr} \ln (E) + \text{Tr} \ln \left\{ 1 + \frac{D^\alpha \bar{D}^2 D_\alpha}{\square^2} M_{AB} \right\} \\ &+ \text{Tr} \ln \left\{ 1 + \xi \left(\frac{\bar{D}^2 D^2}{\square^2} (S_{AB} + T_{AB}) \right. \right. \\ &\left. \left. + \frac{D^2 \bar{D}^2}{\square^2} (S_{BA} + T_{BA}) - \frac{D^2}{\square} R_{AB} - \frac{\bar{D}^2}{\square} \bar{R}_{AB} \right) \right\} . \end{aligned} \quad (4.8)$$

We now want to do the D -algebra in the expression (4.1). To this end we insert 1 in the form $(\{D^2, \bar{D}^2\} - D^\alpha \bar{D}^2 D_\alpha) / \square$ in front of $\text{Tr} \ln \mathcal{O}$ (see Eq. (4.11), below). Using the properties of the projection operators we can convert all the spinor derivatives inside the logarithms into boxes. Let us analyze each term separately.

We extract an irrelevant factor $\propto \text{Tr} \ln(\square)$ from the first term (and drop it in Eq. 4.10 and below):

$$\text{Tr} \ln(E) = \text{Tr} \ln \begin{pmatrix} 0 & \square \\ \square & 0 \end{pmatrix} + \text{Tr} \ln \left\{ \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} + \begin{pmatrix} 0 & \frac{\bar{\mu} D^2}{\square} \\ \frac{\mu \bar{D}^2}{\square} & 0 \end{pmatrix} \right\} . \quad (4.9)$$

Because of the trace, only even powers survive in the expansion of the log and we can rewrite (4.9) as

$$\frac{1}{2} \text{Tr} \ln \left\{ \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} - \begin{pmatrix} \frac{\bar{\mu}\mu D^2 \bar{D}^2}{\square^2} & 0 \\ 0 & \frac{\mu\bar{\mu} \bar{D}^2 D^2}{\square^2} \end{pmatrix} \right\} . \quad (4.10)$$

Performing part of the trace and inserting 1 in the form given above, the final expression is:

$$-\frac{1}{2} \int d^4x d^4\theta d^4\theta' \delta(\theta' - \theta) \text{Tr} \ln(E(x, \theta)) \delta(\theta - \theta') = \quad (4.11)$$

$$-\frac{1}{4} \int d^4x d^4\theta d^4\theta' \delta(\theta' - \theta) \frac{\{D^2, \bar{D}^2\} - D^\alpha \bar{D}^2 D_\alpha}{\square} \times \\ \times \text{Tr} \ln \left(1 - \frac{D^2 \bar{D}^2}{\square^2} \bar{\mu}\mu - \frac{\bar{D}^2 D^2}{\square^2} \mu\bar{\mu} \right) \delta(\theta - \theta') . \quad (4.12)$$

After doing the D -algebra and using the properties of the spinor derivatives and the delta functions, and taking the Fourier transform, this reduces to

$$\frac{1}{2(4\pi)^2} \int d^4\theta dk^2 \text{Tr} \ln \left(1 + \frac{\bar{\mu}\mu}{k^2} \right) , \quad (4.13)$$

where the final trace is over the i, k indices of the $\bar{\mu}\mu$ matrix and this agrees with the scalar multiplet contribution calculated in super-Landau gauge (3.3).

The second term, treated in the same way, becomes

$$-\frac{1}{(4\pi)^2} \int d^4\theta dk^2 \text{Tr} \ln \left(1 + \frac{M}{k^2} \right) , \quad (4.14)$$

where the trace is over the adjoint indices of M_{AB} ; this agrees with our previous (super-Landau gauge) result (3.4).

The third term contains all the ξ dependence and in Landau gauge it vanishes. The D -algebra in this case is not so straightforward as in the previous cases. To perform it we have to use the identity

$$1 + XD^2 \bar{D}^2 + Y \bar{D}^2 D^2 + ZD^2 + \bar{Z} \bar{D}^2 = \\ (1 + \bar{N} \bar{D}^2) (1 + UD^2 \bar{D}^2 + V \bar{D}^2 D^2) (1 + ND^2) , \quad (4.15)$$

where X, Y, Z, \bar{Z} are some arbitrary matrices and

$$\begin{aligned} N &= (1 + X\Box)^{-1} Z \\ \bar{N} &= \bar{Z} (1 + X\Box)^{-1} \\ U &= X \\ V &= Y - \bar{Z} (1 + X\Box)^{-1} Z . \end{aligned} \tag{4.16}$$

The D -algebra implies that a function of D^2 or \bar{D}^2 alone vanishes, and we finally get all of the ξ dependence as a sum of two terms.

$$\begin{aligned} &\frac{1}{2(4\pi)^2} \int d^4\theta dk^2 \left\{ \text{Tr} \ln \left(1 - \xi \frac{1}{k^2} (S + T)_{AB} \right) \right. \\ &\quad \left. + \text{Tr} \ln \left(1 - \xi \frac{1}{k^2} (S + T)_{AB} + \xi^2 \bar{R}_{AC} \left[\frac{1}{k^2 - \xi(S+T)} \right]_{CD} R_{DB} \right) \right\} . \end{aligned} \tag{4.17}$$

We have verified that when Φ is restricted to fields which, after gauge symmetry breaking, are massless and hence on shell at zero momentum, the ξ dependent terms vanish. Unfortunately, it seems impossible to actually evaluate the momentum integral and find an explicit expression for the ξ dependent term in the most general case.

5 Examples

We now consider some examples; we focus on the $N = 2$ case, where our field Φ contains an adjoint scalar ϕ (which we can regard as a matrix) and a number of hypermultiplets Q, \tilde{Q} . First consider the case with one hypermultiplet with mass m and let us look at the ϕ dependence of the effective Kähler potential. The hypermultiplet contribution (4.13) gives two equal factors summing up to

$$\frac{1}{(4\pi)^2} \int dk^2 \text{Tr} \ln \left(1 + \frac{(\bar{m} + \bar{\phi})(m + \phi)}{k^2} \right) \tag{5.1}$$

where the trace is over the representation of the hypermultiplet. The vector multiplet contribution from (4.14) is

$$- \frac{1}{(4\pi)^2} \int dk^2 \text{Tr} \ln \left(1 + \frac{\phi\bar{\phi} + \bar{\phi}\phi}{2k^2} \right) \tag{5.2}$$

where the trace is over the adjoint representation. Since we are only considering the dependence on ϕ , the matrices T and R are zero. However S (see 4.7) is not zero and this gives us a ξ dependent Kähler potential where the ξ dependent piece is

$$\frac{1}{(4\pi)^2} \int dk^2 \text{Tr} \ln \left(1 + \xi \frac{\phi \bar{\phi} - \bar{\phi} \phi}{2k^2} \right) \quad (5.3)$$

We can perform the integral and the trace in each case. The vector multiplet contribution is the same as in [4] but the hypermultiplet contribution changes to

$$K_Q^{fund} = -\frac{1}{(8\pi)^2} \left(s \ln \frac{(\bar{\phi}^2 - \bar{m}^2)(\phi^2 - m^2)}{16\Lambda^4} + t \ln \frac{s+t}{s-t} \right), \quad (5.4)$$

where $s = \bar{m}m + \bar{\phi} \cdot \phi$ and $t = \sqrt{(\bar{m}\phi + m\bar{\phi}) \cdot (\bar{m}\phi + m\bar{\phi}) + (\bar{\phi} \cdot \phi)^2 - \bar{\phi}^2 \phi^2}$. Restricting ϕ to the massless fields, *i.e.*, $[\phi, \bar{\phi}] = 0$, this agrees with the expression one gets by expanding the exact solution of Seiberg and Witten [10].⁵

We find the dependence on the gauge fixing parameter by evaluating the integral and the trace in (5.3). The eigenvalues of the matrix $[\phi, \bar{\phi}]$ are $\left\{ 0, \pm \sqrt{(\bar{\phi} \cdot \phi)^2 - \bar{\phi}^2 \phi^2} \right\}$, so when doing the k^2 integral, we are left with

an *imaginary* contribution to the effective action $i\pi\xi\sqrt{(\bar{\phi} \cdot \phi)^2 - \bar{\phi}^2 \phi^2}$. Such a term in the effective action may appear disturbing, but it is typical of unphysical contributions from longitudinal states that are projected out in Landau gauge.⁶ Indeed, it does not contribute to physical quantities, and, in particular, it vanishes for the massless fields of the theory (when ϕ and $\bar{\phi}$ commute). It also does not arise in the gauge-independent effective action of DeWitt and Vilkovisky [5, 6, 7].

Next we consider N_f hypermultiplets in an arbitrary representation and with $m = 0$, and study the dependence of the Kähler potential on the hypermultiplets. For simplicity we also work in super Landau gauge ($\xi=0$). The

⁵We thank Gordon Chalmers for evaluating the asymptotic expansions of the relevant elliptic integrals.

⁶We thank George Sterman for explaining this to us.

hypermultiplet contribution again is a sum of two equal factors giving

$$\frac{1}{(4\pi)^2} \int dk^2 \text{Tr} \ln \left(\delta_{AB} + \frac{\sum_{a=1}^{N_f} (\bar{Q}^a T_A T_B Q^a + \tilde{Q}^a T_B T_A \tilde{\bar{Q}}^a)}{k^2} \right), \quad (5.5)$$

and the vector multiplet contribution is

$$- \frac{1}{(4\pi)^2} \int dk^2 \text{Tr} \ln \left(\delta_{AB} + \frac{\sum_{a=1}^{N_f} (\bar{Q}^a \{T_A, T_B\} Q^a + \tilde{Q}^a \{T_A, T_B\} \tilde{\bar{Q}}^a)}{2k^2} \right) \quad (5.6)$$

Note that $T_A T_B = \frac{1}{2}(\{T_A, T_B\} + [T_A, T_B]) = \frac{1}{2}(\{T_A, T_B\} + f_{AB}^C T_C)$ implies that the two contributions cancel exactly whenever

$$\sum_{a=1}^{N_f} (\bar{Q}^a T_A Q^a - \tilde{Q}^a T_A \tilde{\bar{Q}}^a) = 0; \quad (5.7)$$

This condition is satisfied by hypermultiplets that remain massless. Thus, massless hypermultiplets receive no corrections to their Kähler potential, as stated in [10]. However, the massive multiplets are expected to receive corrections, and, in particular, we expect to find terms that mix the hypermultiplets and the vector multiplet chiral field ϕ . These are presumably Kähler residues of higher dimension terms analogous to those found for the vector multiplet alone in [4]; however, since a suitable off-shell formulation of the hypermultiplet is not known, we do not even know how to write such terms down.

While we were preparing this manuscript, a preprint [11] appeared on the hep-th archive; the work has some overlap both with this paper and with [4]. The authors do not discuss their gauge choice, but appear to work in supersymmetric Feynman gauge without considering the imaginary term that we found away from Landau gauge.

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